Evidences for pairing of nearly-free quasiparticles from paraconductivity in layered superconducting cuprates

S. Caprara[1], M. Grilli[1], B. Leridon[2] and J. Vanhacken[3]

[1] *INFM-CNR, Unità di Roma 1 and SMC Center, and Dipartimento di Fisica, Università di Roma "Sapienza", piazzale Aldo Moro 5, I-00185 Roma, Italy*
[2] *Laboratoire de Physique Quantique - ESPCI/UPR5-CNRS, 10, Rue Vauquelin - 75005 Paris - France*
[3] *INPAC, Institute for Nanoscale Physics and Chemistry, Katholieke Universiteit Leuven, Celestijnenlaan 200 D, B-3001 Heverlee, Belgium*

PACS 74.72.-h – Cuprate superconductors (high-Tc and insulating parent compounds)
PACS 74.25.Fy – Transport properties (electric and thermal conductivity, thermoelectric effects, etc.)
PACS 74.40.+k – Fluctuations (noise, chaos, nonequilibrium superconductivity, localization, etc.)
PACS 74.20.De – Phenomenological theories (two-fluid, GinzburgLandau, etc.)

Abstract. - We revisit the Aslamazov-Larkin theory of paraconductivity in two dimensions, to distinguish its universal features from the specific features of nearly-free paired fermions. We show that both the numerical prefactor and the temperature dependence of the experimental paraconductivity in underdoped $La_{2-x}Sr_xCuO_4$ are only compatible with pairing of nearly-free fermionic quasiparticles. This conclusion is strengthened by the analysis of paraconductivity data in the presence of a finite magnetic field, from which we extract a rather low value of the critical field $H_{c2}(T=0)$.

Introduction. – Layered superconducting (SC) cuprates are characterized by a pseudogap state in the underdoped region of the phase diagram, below a temperature T^* which at low doping is much larger then the SC critical temperature T_c and merges with it near optimal doping, where T_c is maximum. A possible explanation relies on the formation of incoherent SC Cooper pairs below T^*, the modulus $|\Delta|$ of the SC order parameter acting as the pseudogap detected by various thermodynamical and transport measurements. Superconductivity is prevented by fluctuations of the phase of the order parameter, and develops only below T_c, where phase coherence is eventually established and the preformed pairs condense. The observation of a sizeable Nernst effect [1] and a strong diamagnetic response [2, 3] above T_c have been interpreted in this sense [1]. If this were the case, however, the most anisotropic cuprates [e.g., $Bi_2Sr_2CaCu_2O_{8+\delta}$ (BSCCO)] should display an exponential temperature dependence in the enhancement of conductivity due to SC fluctuations at temperatures $T > T_c$ [the so-called *paraconductivity*], associated with vortical fluctuations, typical of a Kosterlitz-Thouless transition in two dimensions (2D) [4]. Instead, it is well documented that paraconductivity in all the families of cuprates is fully accounted for by the standard Aslamazov-Larkin (AL) theory [5, 6] based on Gaussian SC fluctuations, with the real and imaginary part of the SC order parameter Δ fluctuating around zero. While $YBa_2Cu_3O_{7-x}$ is less anisotropic and displays the AL behavior characteristic of three-dimensional systems [7], all other compounds, which have a more anisotropic structure, display the standard AL behavior for two-dimensional systems (see, e.g., Refs. [8, 9] and references therein). In particular recent experiments in underdoped $La_{2-x}Sr_xCuO_4$ (LSCO) recovered the normal state under strong magnetic field, thereby allowing for an unambiguous determination of paraconductivity [9], *leaving no room for a contribution of vortical phase fluctuations* over the broad temperature range relevant for the pseudogap. This result challenges the phase-fluctuation scenario raising the following issue: How stringent is the above conclusion based on the AL expression for paraconductivity? Within a general phenomenological Ginzburg-

[1]This interpretation was recently questioned, and an alternative explanation in terms of Gaussian SC fluctuations was proposed, see, e.g., L. Cabo, *et al.*, Phys. Rev. Lett. **98**, 119701 (2007); N. P. Ong, *et al.*, *ibid.*, 119702 (2007).

Landau (GL) approach, we show that the AL functional form in 2D [$\propto (T-T_c)^{-1}$] is fairly general because ultimately stems from two general principles, namely gauge-invariance and the hydrodynamic form of the pair collective modes. On the contrary, the numerical prefactor is specific of the fermionic state and therefore provides valuable information on the microscopic state of the system. Specifically, we show here that the precise AL value of the paraconductivity coefficient stems from the assumption of fermions with very narrow spectral weight (i.e., nearly free fermionic quasiparticles). This result is one of the two central points of this Letter and, together with experiments of Ref. [9], which dictate the specific value and temperature dependence of this factor, clearly indicates that in underdoped LSCO fluctuations not only are Gaussian, but also arise from pairing of apparently *weakly-coupled* fermionic quasiparticles.

To challenge this quite surprising result, we present here new paraconductivity data in weak magnetic fields. We find that paraconductivity is still fully compatible with weakly paired quasiparticles and we also introduce the new concept of "hidden" critical field at zero temperature, $H_{c2}^G(0)$, related to the Gaussian fluctuations only. Its value is remarkably lower than the one usually reported in the literature, strengthening our conclusion that paraconductivity is related to superconductivity due to weakly-coupled quasiparticles. This is the second remarkable point of this work. These evidences of weakly-coupled quasiparticles, are surprising because their presence could hardly be guessed from the quite anomalous form of the normal state resistivity and is at odds with the broad spectral lines usually observed in photoemmission experiments in cuprates [10]. Our aim is not to solve this apparent contradiction, but rather to draw attention to this feature. To extract all information from the data, we preliminarly revisit the theoretical derivation of the Gaussian theory, putting precise bounds to the meaning and generality of the 2D AL expression.

Gauge invariant hydrodynamic description of paraconductivity. – A superconductor can be described within a generic model of fermions coupled by a λ-wave pairing interaction (most frequently s- or d-wave have been considered for singlet superconductors). As customary, by integrating out the fermions one derives an effective action for the pair field $\Delta(\mathbf{r},\tau)$ (here $\mathbf{r}$ is the coordinate vector and τ is the imaginary time within the finite-temperature formulation). The quadratic (Gaussian) part of the resulting action is

$$\mathcal{S}_G = \int_0^\beta d\tau \int d^D\mathbf{q}\,\Delta^*(\mathbf{q},\tau)\left[a + C\mathbf{q}^2 + \gamma\partial_\tau\right]\Delta(\mathbf{q},\tau), \tag{1}$$

where D is the space dimensionality, $\Delta(\mathbf{q},\tau)$ is the Fourier transform of $\Delta(\mathbf{r},\tau)$ with respect to $\mathbf{r}$, and $\mathbf{q}$ is the corresponding wavevector. Whereas the explicit expressions of the coefficients a, C, and γ depend on the details of the microscopic model, e.g., the pairing symmetry and the fermionic density of states (DOS), Eq. (1) holds generically whenever a hydrodynamic description for the pair field is adequate, and is indeed phenomenologically adopted in the time-dependent Ginzburg-Landau approach [6]. In this Letter we consider Gaussian fluctuations above T_c, and keep only the action (1), discarding higher-order terms. In the GL approach, one may conventionally take $\gamma = \gamma_{GL} = 1$, which amounts to rescale Δ so that its equation of motion is the Schrödinger equation. Thus, in the Gaussian approximation, physical quantities only depend on two parameters, the *mass* $a_{GL} \equiv a/\gamma$ and the *stiffness* $C_{GL} \equiv C/\gamma$.

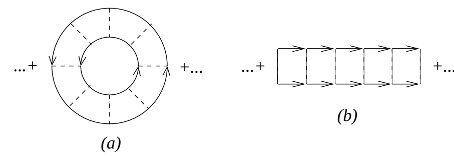

Fig. 1: Typical diagram for the Baym-Kadanoff functional (a) and $\mathcal{T}$-matrix propagator of Gaussian fluctuations (b) adopted in this Letter. Dashed and solid lines represent, respectively, the pairing interaction and the fermion propagator (see text).

The pair field, with a charge $2e$, is coupled to a spatially uniform electromagnetic field $\mathbf{A}(\tau)$ taking $\mathbf{q} \to \mathbf{q} - 2e\mathbf{A}(\tau)$, as dictated by gauge invariance. The AL contribution to the current-current response, and hence to paraconductivity, is associated with the current density $4eC\mathbf{q}\,\Delta^*(\mathbf{q},\tau)\Delta(\mathbf{q},\tau)$, and the prefactor in the current vertex can be identified with the stiffness. Under the assumption of a *gauge-invariant hydrodynamic description* for the SC pair fluctuations, the above arguments hold *irrespective of the Fermi-liquid or non-Fermi-liquid character of the normal state.* Of course, any microscopic derivation of Eq. (1) must obey gauge invariance. In the case of strongly interacting fermions such a derivation is overwhelmingly difficult and beyond the scope of this Letter. On the other hand in the following section, we provide an example of the current-stiffness relation in the case of weakly-coupled fermions.

Weak-coupling microscopic derivation. – Our treatment closely follows the gauge-invariant approach of Ref. [11]. We start from a Baym-Kadanoff functional (i.e., the microscopic equivalent of the GL functional) and obtain the paraconductivity by insertion of current vertices. For weakly-coupled fermions one can adopt the Baym-Kadanoff functional shown in Fig. 1(a).

For definiteness we assume a separable potential $V(\mathbf{k},\mathbf{k}') = Vw_\lambda(\mathbf{k})w_\lambda(\mathbf{k}')$ of strength V, promoting λ-wave pairing [in cuprates, e.g., d-wave, with $w_d = \cos(k_x) - \cos(k_y)$]. A weak-coupling $\mathcal{T}$-matrix approximation yields the pair propagator of Fig. 1(b), i.e., the inverse of the coefficient of the action (1),

$$\mathcal{K}_\lambda(\mathbf{q},\omega_\ell) = \frac{1}{V^{-1} - \Pi_\lambda(\mathbf{q},\omega_\ell)} \approx \frac{1}{a_\lambda + C_\lambda\mathbf{q}^2 + \gamma_\lambda|\omega_\ell|}, \tag{2}$$

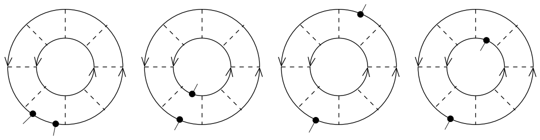

Fig. 2: Diagrams of the current-current response functions generated from the Baym-Kadanoff functional of Fig. 1(a): DOS correction (a), Maki-Thompson vertex correction (b), and AL contributions ($c1$, $c2$). The full circle with a thin line represents a current-vertex insertion (see text).

with the τ variable Fourier-transformed into the Matsubara frequency ω_ℓ. The λ-wave particle-particle bubble is

$$\Pi_\lambda(\mathbf{q},\omega_\ell) \equiv T\sum_{\mathbf{k},\varepsilon_n} w_\lambda^2(\mathbf{k})\,\mathcal{G}(\mathbf{k}+\mathbf{q},\varepsilon_n+\omega_\ell)\mathcal{G}(-\mathbf{k},-\varepsilon_n), \tag{3}$$

$\mathcal{G}(\mathbf{k},\varepsilon_n) \equiv (i\varepsilon_n - \xi_\mathbf{k})^{-1}$ is the fermion propagator, and $\xi_\mathbf{k}$ is the fermion dispersion. An expansion of $\Pi_\lambda(\mathbf{q},\omega_\ell)$ at small $\mathbf{q}$ and ω_ℓ yields, respectively, C_λ and γ_λ. The mass $a_\lambda \equiv V^{-1} - \Pi_\lambda(0,0)$ linearly vanishes at $T = T_c$.

The insertion of two current vertices in the diagrams of Fig. 1(a) yields the current-current correlation functions [11] shown in Fig. 2. The diagrams of Figs. 2($c1$) and 2($c2$) give the AL contributions, once the ladder resummation of Fig. 1(b) is adopted for the pair propagator. These contributions are different from the others, as they vanish if the fermionic loops with one current-vertex are evaluated for zero frequency and momentum of the pair propagators, due to the vector character of the current vertex. The first non-zero contribution to each loop is $\tilde{C}\mathbf{q}$ [5], where $\tilde{C}$ is a constant prefactor. Gauge invariance imposes a definite relation between $\tilde{C}$ and the stiffness C. This relation is enforced by a Ward identity, which can be derived from the Baym-Kadanoff functional, and to first order in the momentum difference $\mathbf{s}$ reads

$$\begin{aligned}&\mathcal{K}_\lambda^{-1}(\mathbf{q}+\mathbf{s},\omega_\ell) - \mathcal{K}_\lambda^{-1}(\mathbf{q},\omega_\ell) = \\ &T\sum_{\mathbf{k},\varepsilon_n} w_\lambda^2(\mathbf{k})\mathcal{G}(\mathbf{k},\varepsilon_n)\mathcal{G}(\mathbf{k},\varepsilon_n)\mathcal{G}(-\mathbf{k}+\mathbf{q},-\varepsilon_n)\,\mathbf{v}_\mathbf{k}\cdot\mathbf{s},\end{aligned} \tag{4}$$

where $\mathbf{v}_\mathbf{k} \equiv \partial_\mathbf{k}\xi_\mathbf{k}$ is the fermion velocity, acting as a current vertex in the fermion loops. The direct calculation in the weak-coupling limit yields indeed $2C = \tilde{C}$, where the factor of 2 stems from the $2e$ charge of the pair field.

What can be inferred from observation of AL paraconductivity. – The identification of the coefficient $\tilde{C}$ of the AL current vertex with the stiffness C is the reason why the AL paraconductivity in 2D assumes an expression which is independent of C. Indeed, the AL current-current response reads [5,6]

$$\begin{aligned}\delta\chi_{AL}(\Omega_n) = 4e^2T\sum_{\omega_\ell}\int d^D\mathbf{q}\,&\frac{1}{a_{GL,\lambda}+C_{GL,\lambda}\mathbf{q}^2+|\omega_\ell|}\\ \times\frac{1}{a_{GL,\lambda}+C_{GL,\lambda}\mathbf{q}^2+|\omega_\ell+\Omega_n|}\,&C_{GL,\lambda}^2\mathbf{q}^2,\end{aligned} \tag{5}$$

where the dependence on γ_λ was eliminated in the GL spirit, introducing the two independent parameters $a_{GL,\lambda} \equiv a_\lambda/\gamma_\lambda$, $C_{GL,\lambda} \equiv C_\lambda/\gamma_\lambda$, as discussed above. In the classical limit the sum over ω_ℓ is dominated by the term $\omega_\ell = 0$. After the analytic continuation $i\Omega_n \to \omega + i0^+$, the AL paraconductivity is found as $[\mathrm{Im}\delta\chi_{AL}(\omega)/\omega]_{\omega\to 0}$. In 2D the change of variables $C_{GL,\lambda}\mathbf{q}^2 \to x$ makes $C_{GL,\lambda}$ disappear, yielding the well-known result [5]

$$\delta\sigma_{AL}(\varepsilon) = \frac{e^2}{2\pi\hbar d}\frac{T_c}{a_{GL,\lambda}} \equiv \frac{e^2}{16\hbar d\varepsilon}, \tag{6}$$

where d is the interlayer distance, translating the 2D result into the paraconductivity of a layered system, and $\varepsilon \equiv \pi a_{GL,\lambda}/(8T_c)$ is the dimensionless mass. Eq. (6) stems from the assumption of a gauge-invariant hydrodynamical description for the Gaussian pair fluctuations, which in 2D imposes the independence from $C_{GL,\lambda}$, and is thus generic for 2D Gaussian fluctuations [2].

Since we aim to extract as much physical content as possible from the fitting of experimental data with Eq. (6), we now detail the specific value of the coefficients in the various physical situations. All information on the microscopic physical properties is contained in ε. As soon as the fermion DOS changes with temperature (e.g., with the opening of a pseudogap) one may wonder how this is reflected in the temperature dependence of a_{GL} for the various pairing regimes. In a BCS model of weakly-coupled fermions, the explicit calculation of the particle-particle bubble Π_λ can be carried out, yielding

$$\begin{aligned}\gamma_\lambda &= -\sum_\mathbf{k} w_\lambda^2(\mathbf{k})\int dz\,A(\mathbf{k},z)A(\mathbf{k},-z)\,\partial_z f(z)\\ a_\lambda &= V^{-1} - \sum_\mathbf{k} w_\lambda^2(\mathbf{k})\int dy\,dz\,A(\mathbf{k},y)A(\mathbf{k},z)\mathcal{R}(y,z)\end{aligned}$$

where $f(z)$ is the Fermi function, $\mathcal{R}(y,z) = [1-f(y)-f(z)]/(y+z)$, and $A(\mathbf{k},z)$ is the fermion spectral function. If this latter is narrower than $\partial_z f(z)$, it can be replaced by $\delta(\xi_\mathbf{k}-z)$. In this case, a symmetry-dependent weighted DOS $\mathcal{N}_\lambda \equiv \sum_\mathbf{k} w_\lambda^2(\mathbf{k})\delta(\xi_\mathbf{k})$ appears, generalizing the standard s-wave expressions of the γ and a coefficients [6]. This factor enters both in γ_λ and a_λ, and disappears in $a_{GL} \propto a_\lambda/\gamma_\lambda$ leaving the paraconductivity unaffected by the T dependence of the DOS. It is important to recognize that this result follows from the narrow spectral density of the fermions entering the Cooper channel, and is no longer valid if the spectral density is broad. This suggests that the absence of any additional temperature dependence is the specific signature of paraconductivity from weakly-paired nearly-free quasiparticles. In any case, the numerical prefactor relating a_{GL} to $(T-T_c)$ is model dependent and the standard result $\varepsilon = \log(T/T_c)$ is a spe-

[2]Actually the functional form of Eq. (6) is even more general since it also holds for Kosterlitz-Thouless phase fluctuations with ε exponentially vanishing at T_c [4,12].

cific signature of the BCS weak-coupling limit. Therefore, a *pure* AL contribution [with the specific $e^2/(16\hbar d)$ prefactor and $\varepsilon = \log(T/T_c)$] is hardly mistaken and is a clear indication of nearly-free fermions being *the only* carriers responsible for paraconductivity via the formation of fluctuating weakly-coupled Cooper pairs, independently of their DOS and its possible temperature dependence. This observation is the crucial point of our theoretical analysis: from the data shown below it will allow us to infer that electrons giving rise to paraconductivity in LSCO behave *as if* they were forming Cooper pairs of weakly coupled nearly-free quasiparticles. We now apply these theoretical conclusions to the data obtained in underdoped LSCO.

Evidence of nearly-free quasiparticle pairing. – The resistance of several LSCO samples at different dopings has been recently measured as a function of T with and without strong magnetic fields H [9]. The complete destruction of the SC state at $H = 47$ T uncovers a highly unusual normal state with a resistivity well reproduced, over an extended temperature range below 200 K, by the superposition of a linear and a logarithmic term $\rho_N(T) \equiv \rho(T, H = 47\,\mathrm{T}) = AT - B\ln(T/T_0)$, which naturally introduces a temperature scale at which a minimum in the resistivity occurs in underdoped cuprates under strong magnetic fields [9, 13]. For a sample with $x = 0.09$ and $T_c = 19.0$ K our fit gives $A = 7.54$ $\mu\Omega$cm/K, $B = 490$ $\mu\Omega$cm, and $T_0 = 80.3$ K. We propose no explanation or hypothesis for this unusual normal state and rather focus on the SC state appearing when H is reduced. Following Ref. [9], we define the paraconductivity as $\delta\sigma(T) \equiv \rho^{-1}(T, H = 0) - \rho_N^{-1}(T)$, and report the results in Fig. 3 (black dots) as a function of $\varepsilon \equiv \ln(T/T_c)$, in comparison with the 2D AL result in the BCS limit (solid line). Despite the unusual ρ_N, $\delta\sigma(T)$ is very well described by the standard AL expression *with the pure BCS coefficients, without fitting parameters*. Most importantly, we find that not only the temperature dependence is clearly linear in ε^{-1}, but even the numerical prefactor is that of the weak-coupling theory for nearly-free fermions, within error bars of less than 5%. Since the paraconductivity diverges at T_c, uncertainties in the determination of ρ_N are rather immaterial for $T \approx T_c$ and our finding is quite robust. The contribution of Gaussian fluctuations to paraconductivity spreads over a broad temperature range, $T - T_c \sim T_c$, similarly to what found in underdoped BSCCO [8], where however the need to guess the reference normal state made the analysis much less stringent.

Rewriting $\varepsilon = (\xi_0/\xi)^2$, and assuming $\xi_0 \sim 20$ Å, we can estimate the coherence length ξ of the Gaussian fluctuations. Even for $\varepsilon \approx 0.01$, i.e., $T \approx 1.01T_c$, we find $\xi \sim 10\xi_0 \sim 200$ Å, which is much smaller than the value estimated for Kosterlitz-Thouless vortical phase fluctuations in magnetometry experiments in BSCCO [2]. This discrepancy can hardly be due to the different materials, because paraconductivity experiments in BSCCO [8] give values of ξ consistent with those obtained here for LSCO.

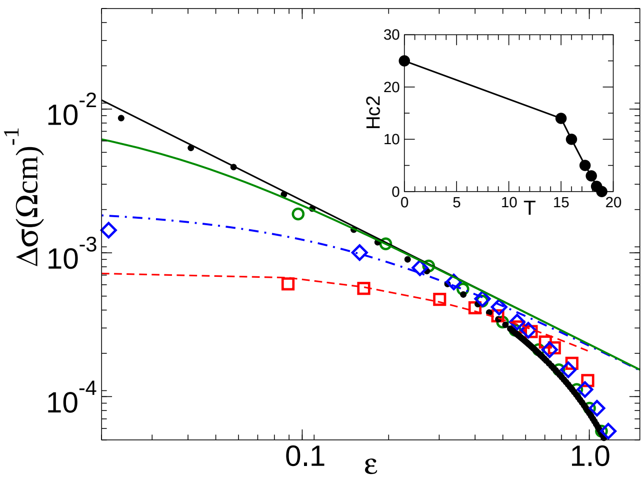

Fig. 3: (Color online) Comparison between the theoretical Gaussian paraconductivity $\delta\sigma(T, H)$ [6] (lines) and the experimental data (symbols) taking an interlayer distance $d = 6.6$ Å [data at $H = 0$ T (black dots), should be compared with theoretical result, Eq. (6), (black solid line)]. $H = 1$ T (green solid line and circles), 5 T (blue dot-dashed line and diamonds), and 14 T (red dashed line and squares). Inset: Gaussian critical temperatures vs H and estimated $H_{c2}(T = 0)$ (see text).

We now focus on new data showing the gradual suppression of Gaussian fluctuations for small-to-moderate H. Since dissipating vortices, introduced by the magnetic field, largely contribute to the resistivity, the Gaussian paraconductivity is difficult to extract. Nevertheless we tested the 2D AL theory at finite H using the expression reported in Ref. [6]. This attempt is obviously meaningful only if the critical temperature *in the absence of dissipating vortices*, $T_c^G(H)$, does not fall deeply into the vortex-dissipation regime. In Fig. 3 we report our results. The choice of $H_{c2}^G(T = 0)$ and of $T_c^G(H)$ is made to optimize the agreement with the data. For $H = 1, 5, 14$ T we find $T_c^G = 18.4, 17.3, 15.0$ K, respectively (see the inset in Fig. 3), which are substantially larger than the experimental $T_c(H)$, determined by vortex dissipation. Therefore our analysis reliably indicates that 2D Gaussian fluctuations persist under substantial magnetic fields. We find $H_{c2}^G(T = 0) = 25$ T, which is much lower than the values at which superconductivity is actually destroyed and usually reported for LSCO at $x = 0.09$ [14]. However, this value is estimated from the weak-coupling expression of Ref. [6], and therefore should be interpreted as *the critical field of a system in which the physics of vortices is absent and only Gaussian fluctuations play a role*. The success of our fitting procedure is based on the existance of a regime where paraconductivity is due to Gaussian fluctuations and would completely fail if only preformed pairs with vortical excitations were present.

Conclusions. – In this work we started from the preliminary remark that AL paraconductivity is ubiquitously

observed in cuprates. This lead us to reexamine the theoretical grounds of AL theory in order to fully ascertain the physical implication of this phenomenological remark. We showed that under general conditions (i.e., gauge invariance and hydrodynamics) 2D paraconductivity is independent of the fluctuation stiffness, and depends on a single parameter, the dimensionless mass ε, which contains all informations on the specific character of the paired fermions. Therefore the robustness of the AL functional form in 2D stems from general physical principles, but the specific numerical prefators may shed light on the nature of the paired fermions. In particular we showed that paraconductivity of the AL functional form with the precise and specific AL prefactors can only be due to weakly-bound nearly-free fermions.

As far as the experimental part of our work is concerned, we concentrated on LSCO only because the new data in strong magnetic field allowed for the unambiguous determination of the reference normal state, but our analysis applies to all families of cuprates. Thus we investigated the experimental paraconductivity in underdoped LSCO showing that it is fully accounted for by Gaussian fluctuations, both in the absence and in the presence of a magnetic field. The supporting theoretical analysis allows to conclude i) that within the experimental errors, in paraconductivity there is no room for contributions due to vortical phase fluctuations, which seem instead to be present in other experimental quantities [1–3]. Moreover ii) the specific value of the numerical prefactor and the temperature dependence of the experimental dimensionless mass indicate that, despite the very anomalous normal state uncovered by the magnetic field, Gaussian fluctuations arise from the pairing of nearly-free fermionic quasiparticles. This would agree with the recent observation of a (small) Fermi surface of nearly-free electrons in underdoped $YBa_2Cu_3O_{6.5}$ [15]. This indication of pairing of weakly coupled quasiparticles, whose presence can hardly be guessed from other physical properties of the cuprates, is perhaps the most surprising and intriguing result of our analysis.

One might speculate that the weakly bound pairs probed by paraconductivity coexist with more tightly bound pairs related to the vortical phenomenology. This coexistance, already implicit in a previous analysis of a two-gap model [16], could also be consistent with recent observations of different gap scales [17].

In this work we also carry out, both experimentally and theoretically, a new analysis by investigating paraconductivity in magnetic field, introducing the new concept of "hidden" critical field related to weakly-bound pairs, which is usually masked by the vortex physics and which rules the destruction of Gaussian fluctuations around T_c.

* * *

We are indebted with C. Castellani, C. Di Castro, J. Lesueur and A. Varlamov for interesting discussions. S.C. and M.G. acknowledge financial support from MIUR-PRIN 2005 - prot. 2005022492. S.C., M.G. and B.L. acknowledge support from CNRS PICS #3368. B.L. also acknowledges the ESF for support through the THIOX short visit grant number 1081.

REFERENCES

[1] Ong N. P., Wang Y., Ono S., Ando Y. and Uchida S., *Ann. Phys.*, **13** (2004) 9, and references therein.

[2] Lu Li, Wang Y., Naughton M. J., Ono S., Ando Y. and Ong N. P., *Europhys. Lett.*, **72** (2005) 451.

[3] Wang Y., Lu Li, Naughton M. J., Gu G. D., Uchida S. and Ong N. P., *Phys. Rev. Lett.*, **95** (2005) 247002.

[4] Halperin B. I. and Nelson D. R., *J. Low. Temp. Phys.*, **36** (1979) 599.

[5] Aslamazov L. G. and Larkin A. I., *Phys. Lett. A*, **26** (1968) 238; *Sov. Phys. Solid State*, **10** (1968) 875.

[6] Larkin A. and Varlamov A., *Theory of fluctuations in superconductors* (Clarendon Press, Oxford) 2005.

[7] Leridon B., Défossez A. , Dumont J., Lesueur J. and Contour J. P., *Phys. Rev. Lett.*, **87** (2001) 197007 and references therein.

[8] S. Caprara, Grilli M., Leridon B., and Lesueur J., *Phys. Rev. B*, **72** (2005) 104509.

[9] Leridon B., Vanacken J., Wambecq T. and Moshchalkov V. V., *Phys. Rev. B*, **76** (2007) 012503.

[10] Damascelli A., Hussain Z. and Shen Z.-X., *Rev. Mod. Phys.*, **75** (2003) 473.

[11] V. Ambegaokar, *Superconductivity*, edited by P.R. Wallace (Gordon and Breach, New York) 1969.

[12] Sachdev S., *Physica C*, **408-410** (2004) 218.

[13] Boebinger G. S., Ando Yoichi, Passner A., Kimura T., Okuya M., Shimoyama J., Kishio K., Tamasaku K., Ichikawa N. and Uchida S., *Phys. Rev. Lett.*, **77** (1996) 5417; Ono S., Ando Yoichi, Murayama T., Balakirev F. F., Betts J. B. and Boebinger G. S., *Phys. Rev. Lett.*, **85** (2000) 638.

[14] Lu Li, CHECKELSKY J. G., KOMIYA SEIKI, ANDO YOICHI and ONG N. P., *Nature Physics*, **3** (2007) 311.

[15] Doiron-Leyraud N., Proust C., LeBoeuf D., Levallois J., Bonnemaison J.-B., Liang R., Bonn D. A., Hardy W. N. and L. Taillefer, *Nature*, **447** (2007) 565.

[16] Perali A., Castellani C., Di Castro C., Grilli M., Piegari E. and Varlamov A. A., *Phys. Rev. B*, **62** (2000) R9295.

[17] Le Tacon M., Sacuto A., Georges A., Kotliar G., Gallais Y.,Colson D. and Forget A., *Nature Phys.*, **2** (2006) 537; Tanaka K., Lee W. S., Lu D. H., Fujimori A., Fujii T., Risdiana, Terasaki I. , Scalapino D. J., Devereaux T. P. and Hussain Z., *Science*, **314** (2006) 1910.